

Constraints on the emission mechanisms of gamma-ray bursts

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ABSTRACT

If the emission of gamma-ray bursts were due to the synchrotron process in the standard internal shock scenario, then the typical observed spectrum should have a slope $F_\nu \propto \nu^{-1/2}$, which strongly conflicts with the much harder spectra observed. This directly follows from the cooling time being much shorter than the dynamical time. Particle re-acceleration, deviations from equipartition, fastly changing magnetic fields and adiabatic losses are found to be inadequate to account for this discrepancy. We also find that in the internal shock scenario the relativistic inverse Compton scattering is always as important as the synchrotron process, and faces the same problems. This indicates that the burst emission is not produced by relativistic electrons emitting synchrotron and inverse Compton radiation.

Key words: gamma rays: bursts — X-rays: general — radiation mechanisms: non-thermal

1 INTRODUCTION

Since the observational breakthrough by *BeppoSAX* (Costa et al. 1997; van Paradijs et al. 1997) the physics of gamma-ray bursts (GRB) has started to be disclosed. The huge energy and power releases required by their cosmological distances support the fireball scenario (Cavallo & Rees 1978; Rees & Mészáros 1992; Mészáros & Rees 1993), whose evolution and behavior is (unfortunately) largely independent of their origin.

We do not know yet in any detail how the GRB event is related to the afterglow emission, but in the most accepted picture of formation of and emission from internal/external shocks (Rees & Mészáros 1992; Rees & Mészáros 1994; Sari & Piran 1997), the former is due to collisions of pairs of relativistic shells (internal shocks), while the latter is generated by the collisionless shocks produced by shells interacting with the interstellar medium (external shocks). The short spikes ($t_{\text{var}} \sim 10$ ms) observed in the high energy light curves suggest that shell-shell collisions occur at distances $R \simeq 10^{12} - 10^{13}$ cm from the central source, involving plasma moving with bulk Lorentz factor $\Gamma \geq 100$. The fireball starts to be decelerated by the interstellar medium further out, at a distance which depends on the density of this material.

The main radiation mechanism assumed to be responsible for both the burst event and the afterglow is synchrotron (Rees & Mészáros 1994; Sari, Narayan & Piran 1996; Sari & Piran 1997; Panaitescu & Mészáros 1998 – see

however Thompson 1994; Liang 1997; Ghisellini & Celotti 1999; Celotti & Ghisellini 1999; Stern 1999). This requires acceleration of electrons up to ultra-relativistic energies and the presence of a significant magnetic field. Evidence supporting that the afterglow emission is due to the synchrotron process include the power law decay in time of the afterglow flux (for reviews see Piran 1999; Mészáros 1999) and the recently detected linear polarization in GRB 990510 (Covino et al. 1999; Wijers et al. 1999), but the only piece of circumstantial evidence in favor of a synchrotron origin of the burst radiation comes from the predicted frequency of the peak of the burst spectrum. Indeed, it is remarkable that the simple assumption of equipartition among protons, electrons and magnetic field energy densities leads – in the internal shock scenario (ISS) – to a typical emission frequency in agreement with observations.

However, the very same ISS inevitably predicts very fast radiative cooling of the emitting particles. In this Letter we point out that this implies an emitted spectrum much steeper than observed. Although other authors have already pointed out that, in the presence of radiative losses, the predicted spectrum is steep (Cohen et al. 1997; Sari, Piran & Narayan 1998; Chiang 1999), here possible alternatives to avoid this conclusion, in the context of the ISS, are discussed, but found inadequate to account for the discrepancy. In addition, we examine the role of the relativistic inverse Compton process in the ISS, which results to be as impor-

tant as the synchrotron one. Therefore in this scenario the high energy radiation would always be energetically significant, thus requiring a careful estimate of the importance of photon–photon collisions leading to electron–positron pair production.

2 THE ‘STANDARD’ SYNCHROTRON SCENARIO

Let us briefly summarize the main features of the ISS (for simplicity we will also refer to it as the ‘standard’ model). The emission of the burst, which originates from the conversion of bulk kinetic energy into random energy, has a duration which is determined by the central engine. In particular, in order to generate intermittent and complex variability patterns, the engine has to produce/eject several shells of matter. If these propagate with different Lorentz factors, a faster shell will catch up with a slower one and in the interaction a shock will develop, which is assumed to be responsible for the acceleration of electrons to ultra-relativistic energies. These would then lose energy by radiating synchrotron photons.

2.1 Typical radii

If two shells move with different Lorentz factors, Γ and $a\Gamma$ (with $a > 1$), and are initially separated by R_o , they interact at a distance R_i from the central engine, where

$$R_i \simeq \frac{2a^2}{a^2 - 1} R_o \Gamma^2. \quad (1)$$

As a reference value in the following we will adopt $a = 2$, corresponding to $R_i \simeq (8/3)R_o\Gamma^2$. The initial temporal separation of the two shells, $\sim R_o/c$, also determines the duration of the emission produced by a single shell–shell collision, as measured by an external observer.

For $a = 2$ the final bulk Lorentz factor of both shells is $\sim 1.41\Gamma$ and if they have equal masses the difference between the initial and final energy is ~ 6 per cent of their total kinetic energy. In the ISS, this energy is shared among magnetic field, protons and electrons.

The other typical distance characterizing the fireball evolution is the transparency radius R_t , at which an expanding shell becomes optically thin to Thomson scattering. Assuming that each shell carries an energy $E_s = 10^{50} E_{s,50}$ erg in bulk motion, $R_t \simeq 6 \times 10^{12} (E_{s,50}/\Gamma^2)^{1/2} \text{ cm}^*$.

The standard scenario requires $R_i > R_t$, i.e. $\Gamma \gtrsim 350 (E_{s,50}/R_{0,7}^2)^{1/5}$, in order for the radiation produced to freely escape (see e.g. Lazzati, Ghisellini & Celotti 1999). In the following E_s and L_s stand for the kinetic energy and power of each shell, respectively, while L stands for the observed luminosity.

2.2 The typical synchrotron frequency

During the shell interaction electrons are instantaneously accelerated in a collisionless shock, and reach a random

Lorentz factor which corresponds to equipartition with the other forms of energy, i.e. $\gamma_{\text{eq}} = (\Gamma' - 1)n_p m_p / (n_e m_e)$, where Γ' is the Lorentz factor of one shell in the rest frame of the other, and n_p and n_e are the proton and lepton densities, respectively. In the ISS, these are assumed equal (i.e. electron–positron pairs do not significantly contribute to n_e). Deviations from the equipartition value are parametrized by a dimensionless coefficient $\epsilon_e = \gamma/\gamma_{\text{eq}}$.

The out-flowing plasma is magnetized, and a typical/indicative field value is estimated by assuming that either a significant fraction of the power is carried as Poynting flux or the field energy in the emitting region constitutes some fraction ϵ_B of the randomized energy. Both possibilities imply that at the distance where the shells interact the Poynting flux carries a power $L_B \equiv R^2 \Gamma^2 B'^2 c/2 = \epsilon_B L_s$, where $L_s = 4\pi R^2 \Gamma^2 n'_p m_p c^3$ is the kinetic power carried by a single shell. This corresponds to

$$B'_{\text{eq}} = [8\pi \epsilon_B n'_p m_p c^2]^{1/2} = \left(\frac{2\epsilon_B L_s}{c} \right)^{1/2} \frac{1}{\Gamma R} \quad (2)$$

From the above estimates it follows that the typical observed synchrotron frequency is $\nu_{\text{peak}} = 2e/(3\pi m_e c) \gamma^2 B' \Gamma / (1+z)$, which at R_i gives

$$h\nu_{\text{peak}} \simeq 4 \frac{\epsilon_e^2 (\Gamma' - 1)^2 \epsilon_B^{1/2} L_{s,52}^{1/2}}{R_{i,13} (1+z)} \text{ MeV}. \quad (3)$$

As mentioned in the Introduction, the success of the standard scenario in (simply) predicting a typical observed frequency in remarkable agreement with observations is probably the strongest piece of evidence pointing towards the synchrotron process as responsible for the burst emission.

Note that the ‘equipartition coefficients’, ϵ_B and ϵ_e , must be close to unity[†] for the observed value of ν_{peak} to be recovered (note that $(\Gamma' - 1)$ is of order unity in the ISS). In turn this also implies/requires that electron–positron pairs cannot significantly contribute to the lepton density.

The predicted synchrotron spectrum, produced by a quasi mono-energetic particle distribution, has a flux density $F_\nu \propto \nu^{1/3}$ up to the cutoff frequency ν_{peak} . The average observed spectra in the hard X-ray band, which we stress are typically derived from ~ 1 s integrated fluxes, are not inconsistent with this shape. Nevertheless exceptions exist, including spectra much flatter than $\nu^{1/3}$, which have already cast some doubts on the synchrotron scenario (Preece et al. 1998; Lloyd & Petrosian 1999).

In the following we point out that, just because the integration and the dynamical timescales are much longer than the particle cooling timescales, *the expected synchrotron spectrum in the entire X- and soft γ -ray band should have a slope $F_\nu \propto \nu^{-1/2}$* . This dramatically exacerbates the discrepancy between the predictions of the standard scenario and observations.

* Here and in the following we parametrize a quantity Q as $Q = 10^x Q_x$ and adopt cgs units. Primed quantities are evaluated in the comoving frame.

[†] The only estimates of these parameters are inferred from the adoption of this model for the interpretation of the afterglow emission. This suggests values of ϵ_B and ϵ_e substantially smaller than unity.

3 RADIATIVE COOLING TIME AND TIME INTEGRATED SPECTRUM

Consider the radiative cooling timescale (in the observer frame) of typical particles emitting synchrotron (and self-Compton) radiation within the frame of the ISS:

$$\begin{aligned} t_{\text{cool}} &= \frac{\gamma}{\dot{\gamma}} \frac{1+z}{\Gamma} = \frac{6\pi m_e c(1+z)}{\sigma_T B^2 \Gamma \gamma (1+U_r/U_B)} \\ &= 1.14 \times 10^{-7} \frac{\epsilon_e^3 (\Gamma' - 1)^3 \Gamma_2}{\nu_{\text{MeV}}^2 (1+U_r/U_B)(1+z)} \text{ s}, \end{aligned} \quad (4)$$

where U_r and U_B represent the radiation and magnetic energy densities, respectively. As already mentioned, the shortest integration times are of the order of 1 s: this implies that the observed spectrum is produced by a cooling particle distribution. Note also that the cooling timescale is much shorter than the dynamical one, resulting in a relatively efficient radiative dissipation: in this situation adiabatic energy losses are therefore negligible (Cohen et al. 1997; see below). In particular, after one dynamical time $t_d = 10^{-2} t_{d,-2}$ s, the cooling electrons emit at the (cooling) observed frequency $\nu_{\text{cool}} \sim 2.2 \times 10^{14} (1+z) t_{d,-2}^{-2} \Gamma_2^{-1} B_4^{-3} (1+U_r/U_B)^{-2}$ Hz, independent of Γ' . Invoking smaller values of the magnetic field to slow down the cooling and obtain ν_{cool} of the order of few hundreds keV does not help, since in this case the self-Compton emission dominates over the synchrotron one (see below).

Since $t_{\text{cool}} \propto 1/\gamma$, in order to conserve the particle number, the instantaneous cooling distribution has to satisfy $N(\gamma, t) \propto 1/\gamma$. When integrated over time, the contribution from particles with different Lorentz factors is ‘weighted’ by their cooling timescale $\propto 1/\gamma$. Therefore the predicted (time integrated) flux spectrum between ν_{cool} and ν_{peak} is (e.g. Piran 1999)

$$F_\nu \propto t_{\text{cool}} N(\gamma) \dot{\gamma} \frac{d\gamma}{d\nu} \propto \frac{\gamma}{\dot{\gamma}} \frac{1}{\gamma} \dot{\gamma} \nu^{-1/2} \propto \nu^{-1/2}, \quad (5)$$

extending from $\sim h\nu_{\text{cool}} \sim \text{eV}$ to $h\nu_{\text{peak}} \sim \text{MeV}$ energies. We thus conclude that, within the assumptions of the ISS, a major problem arises in interpreting the observed spectra as synchrotron radiation. Let us consider in turn alternative hypotheses, within the same general frame, which might ameliorate this difficulty.

3.1 Deviations from equipartition?

If one maintains the requirement of observing synchrotron photons at $\sim \text{MeV}$ energies, the radiative cooling timescale is almost independent of ϵ_B , i.e. of the assumed value of the magnetic field (the only dependence being through the ratio U_r/U_B ; eq. 4). Furthermore – again from eq (4) – $t_{\text{cool}} \sim t_d$ requires a value of ϵ_e close to ~ 40 , thus violating energy conservation (as the electrons would have more energy than the available one).

On the other hand, if the condition that $\nu_{\text{peak}} \sim 1$ MeV is produced by synchrotron is relaxed, the magnetic field intensity can be smaller than the equipartition value with a consequently longer synchrotron cooling timescale. However, as the radiation energy density has to be of the order of $U_r' \sim L_s/(4\pi R^2 \Gamma^2 c) \sim U_{B,\text{eq}}'$ to account for the observed fluxes, the inverse Compton cooling would be in any case extremely efficient, leading again to short cooling timescales. Therefore

even if the observed radiation is produced by self-Compton emission of relativistic particles, we face the same problem of cooling timescales being so short that the spectrum would be steep, as discussed below.

3.2 Particle re-acceleration?

A further possibility to escape the above conclusion is to assume that particles are continuously re-heated, thus avoiding the formation of a cooled particle distribution. However in the standard ISS new particles are continuously swept by the shock and are all accelerated to the equipartition energy. This is a crucial assumption in order to produce a typical observed peak frequency around a few hundred keV. It is thus not possible – in this scenario – to continuously re-accelerate the very same particles, as the energy required would exceed the available one.

Alternatively, relaxing the requirement of the standard scenario, one can envisage a situation in which only ‘selected’ particles are steadily accelerated for the entire duration of the shell-shell interaction. In this case an extreme fine (and unlikely) tuning is required: in fact, to be consistent with the total energetics, the selected particles have to: i) be fixed in number (only a fraction $\sim t'_{\text{cool}}/(\Delta R'/c)$ of the total number of particles can be accelerated); ii) be always the same; iii) achieve $\gamma \sim \gamma_{\text{eq}}$ even in the absence of an equipartition argument.

It would be also plausible to assume that the emission is produced by a power-law distribution of electrons resulting from continuous acceleration and cooling. Indeed, for an energy distribution $\propto \gamma^{-p}$ (with $p > 0$) only a minority of particles attain the maximum energy. But – besides having to keep all the particles accelerated for the entire duration of the shell-shell interaction – the relative number of the most energetic particles requires $p > 2$, leading to a spectrum even steeper than $\nu^{-1/2}$.

We therefore conclude that re-acceleration does not avoid the spectral discrepancy, even when relaxing some of the key assumptions of the standard scenario.

3.3 Strongly varying magnetic field?

Let us consider the case in which the magnetic field attains a value close to the equipartition one only in a very limited region (e.g. near the shock front), while is weaker elsewhere. In this situation the synchrotron cooling is mostly effective within this region only and the particles may not have time to significantly cool. Therefore in principle a synchrotron spectrum $\propto \nu^{1/3}$ might be produced. This requires that particles loose much less than half of their energy in the radiative zone, since even a reduction of a factor two of their Lorentz factor would imply that a spectrum $\nu^{-1/2}$ is produced in a range spanning a factor four in frequency (this may correspond to the entire BATSE energy range).

Therefore it would be required that:

- i) the synchrotron process in the most radiative region *must* be inefficient, since it has to reduce the electron energy at most by a small fraction;
- ii) away from this zone, particles continue to rapidly cool by self-Compton and – at a reduced rate – by synchrotron emission. The inverse Compton process then becomes the dominant cooling mechanism.

iii) since the cooling is very rapid anyway, the self-Compton emission itself would produce a time integrated (over t_d) steep spectrum.

We conclude that the net effect of having a strong magnetic field confined in a limited region is to decrease the total synchrotron power in favor of the self-Compton one, whose spectrum would in any case be steep (see below).

3.4 Adiabatic losses?

Suppose that particles are accelerated in compact regions that rapidly expand because of internal pressure. Adiabatic losses dominate particle cooling as soon as the particle Lorentz factor decreases below some critical γ_{ad} , thus generating a spectrum $\propto \nu^{1/3}$ below the synchrotron frequency ν_{ad} (corresponding to γ_{ad}), and steeper above (Cohen et al. 1997). However this possibility faces two severe problems, both related to the overall efficiency, being required that:

- i) each electron loses only a small fraction of its energy radiatively (i.e. $\gamma_{\text{ad}}/\gamma$ must be greater than $\sim 1/2$);
- ii) the emitting regions are very compact, for adiabatic losses to be significant. This implies that the transformation of bulk into random energy does not occur in a shell subtending the entire ejection solid angle. Photon and electron densities have then to be higher to account for the observed luminosity, thus enhancing the inverse Compton process.

4 IMPORTANCE OF THE RELATIVISTIC INVERSE COMPTON PROCESS

As long as the scattering optical depth τ_T of the electron in the emitting region is smaller than unity, the importance of the relativistic inverse Compton process with respect to the synchrotron one is measured by the (relativistic) Compton parameter $y' \equiv \tau_T \gamma^2 \beta^2 = \sigma_T \gamma^2 \beta^2 n'_e ct'_{\text{cool}}$. The width of the region corresponds to a cooling length, ct'_{cool} , as assumed within the standard ISS.

Since the magnetic field intensity is related to the proton density n'_p we obtain

$$y' = \frac{3}{4} \frac{\epsilon_e}{\epsilon_B} \frac{n'_e}{n'_p} \frac{\Gamma' - 1}{1 + U'_r/U'_B}. \quad (6)$$

This implies that the inverse self-Compton power is of the same order of the synchrotron one ‡ , and is emitted at a typical observed energy (for the first order)

$$h\nu_c \simeq \gamma^2 h\nu_s \simeq 13 \frac{\epsilon_e^4 (\Gamma' - 1)^4 \epsilon_B^{1/2} L_{s,52}^{1/2}}{R_{i,13} (1+z)} \text{ TeV}. \quad (7)$$

Two points are worth being stressed. First, the strong dependence of ν_c on the equipartition parameter ϵ_e . Furthermore, although in the comoving frame this typical Compton frequency is a factor Γ lower, it can still largely exceed the pair production threshold (see below).

It has been mentioned and implicitly assumed above that even in the hypothesis that the hard X-ray burst radiation is due to self-Compton emission, the argument of the fast cooling producing a steep spectrum applies. Let us

‡ This is true as long as the scattering process is in the Thomson regime, i.e. $\gamma < 800 B_4'^{-1/3}$ for the first order Compton scattering.

consider this possibility more closely, and in particular the first order inverse Compton spectrum.

Although in this case the typical electron energies required are smaller, the cooling timescales are still much shorter than the dynamical time: in fact, to produce $\sim \text{MeV}$ photons by the first order Compton scattering, $\gamma \sim 83 [\nu_{\text{MeV}} (1+z)/(\Gamma_2 B_4')]^{1/4}$ with a corresponding cooling time

$$t_{\text{cool}} = 1.4 \times 10^{-5} \frac{R_{i,13}^2 \Gamma_2^{5/4} B_4'^{1/4} (1+z)^{3/4}}{L_{50} (1 + U_B/U_r) \nu_{\text{MeV}}^{1/4}} \text{ s}. \quad (8)$$

Here $L = 10^{50} L_{50} \text{ erg s}^{-1}$ is the (observed) luminosity produced by a single shell.

Furthermore, by following the same arguments leading to eq. (5), the predicted time-integrated spectrum results $F_\nu \propto \nu^{-3/4}$, i.e. *even steeper* than $\nu^{-1/2}$.

A further difficulty of interpreting the burst emission as first order scattering is that if the inverse Compton power exceeds the synchrotron one by a certain factor, then each higher Compton order will dominate over the previous one by the same amount, until the typical emitted frequency reaches the electron energy. Only a small fraction of the radiated power would therefore be observed (in the hard X-ray band).

5 PAIR PRODUCTION

The above results indicate that the time integrated spectrum predicted by the standard scenario is steeper than observed. Furthermore, the power emitted through the self-Compton process should be comparable to – if not more than – the synchrotron one, and emitted at energies exceeding the pair production threshold. It is thus compelling to estimate the importance of photon-photon collisions producing electron-positron pairs.

Setting $x \equiv h\nu/(m_e c^2)$, the energy threshold for photons of energy x is $x_T = 2/[x(1 - \cos \theta)]$, where θ is angle between the two photon directions. Also, the photon-photon collision rate is proportional to $(1 - \cos \theta)$.

The result of the integration of the photon-photon cross section over the energy of the target photons can be well approximated by $(\sigma_T/5) x_T n_\gamma(x_T)$ (Svensson 1987), where $n_\gamma(x_T)$ is the number density of photons of energy x_T , which is related to the observed luminosity $L(x_T)$ by $x_T n_\gamma(x_T) = L(x_T)/(4\pi m_e c^3 R^2)$.

The optical depth for pair production in the observer frame can be then expressed as

$$\tau_{\gamma\gamma}(x) = \frac{\sigma_T}{20\pi} \frac{L(x_T) \langle 1 - \cos \theta \rangle}{R m_e c^3} \frac{\Delta R}{R}. \quad (9)$$

ΔR may represent the width of the emitting shell or, alternatively, the typical scale over which the emitted photons might interact (i.e. $\Delta R \sim R$), depending whether we are interested in the pair production within the shell or also outside it.

Since the source is moving relativistically, all photons appear to be emitted quasi-radially, and interact with a typical angle $\sin \theta \sim 1/\Gamma$ (corresponding to $\cos \theta \sim \beta$), and thus $\langle 1 - \cos \theta \rangle \sim 1/\Gamma^2$. If the typical size of the fireball is estimated by time variability, $R \sim ct_{\text{var}} \Gamma^2$, we have:

$$\tau_{\gamma\gamma}(x) = \frac{\sigma_T}{20\pi} \frac{L(x_T)}{t_{\text{var}} m_e c^4 \Gamma^4} \frac{\Delta R}{R} \quad (10)$$

In this form $\tau_{\gamma\gamma}(x)$ can be estimated even without detailed spectral information. If the observed spectrum is a power law $L(x) \propto x^{-\alpha}$, with $\alpha < 1$ up to a maximum energy x_{\max} , the observed luminosity at threshold is related to the total luminosity L by $L(x_T) = [(1-\alpha)/2^\alpha](L/x_{\max}^{1-\alpha})(x/\Gamma^2)^\alpha$, giving

$$\tau_{\gamma\gamma}(x) = \frac{(x/2)^\alpha}{\Gamma^{4+2\alpha}} \frac{(1-\alpha)\ell}{20\pi x_{\max}^{1-\alpha}} \frac{\Delta R}{R}, \quad (11)$$

where the compactness $\ell \equiv \sigma_T L / (t_{\text{var}} m_e c^4)$ has been introduced. Note that the optical depth increases with photon energy x .

For illustration, consider a burst with $\Gamma = 10^2$ lasting $t_{\text{var}} = 10$ ms. 1 GeV photons ($x = 2000$) mostly interact with target photons of energies $x_T = 10 \Gamma_2^2$, i.e. ~ 5 MeV. Assume that the observed luminosity at these energies is $L(x_T) = 10^{50}$ erg s $^{-1}$. From eq. (10) we have

$$\tau_{\gamma\gamma}(x = 2000) \sim 1.4 \times 10^3 \frac{L_{50}(x_T)}{\Gamma_2^4} \frac{\Delta R}{R}. \quad (12)$$

Since the optical depth is so large, all the high (\sim GeV) energy emission can be easily absorbed, unless $\Gamma > 10^3$.

For $\Gamma < 10^3$ relativistic pairs can then be copiously produced. They will immediately cool radiatively, initiating a pair cascade, strongly affecting the primary spectrum, and possibly even the dynamics. (Equilibrium) spectra produced by pair cascades have been extensively studied in the past in the context of nuclear AGN emission. The general outcome is that pairs act as reprocessors of the high energy emission, which is absorbed and ultimately reprocessed into lower energy radiation. This corresponds to a steepening of the spectrum, thus exacerbating the discrepancy with the observed bursts.

Furthermore, the lepton density n_e may become substantially larger than the proton one, n_p , and thus the equipartition energy γ_{eq} smaller than m_p/m_e .

6 DISCUSSION

The main point stressed in this paper is the inadequacy of the synchrotron and inverse Compton emissions from ultra-relativistic electrons to account for the observed burst spectra – at least within the scenario invoking internal shocks for the dissipation of the fireball bulk kinetic energy.

This is independent of detailed assumptions and directly follows from the extremely short cooling timescales – compared to the dynamical time which is in turn much smaller than the integration one – required by the ISS, which lead to a steep emitted spectrum. No alternative hypothesis, which could alleviate this spectral discrepancy, has been found. Furthermore it is stressed that within the ISS scenario electron–positron pairs would be naturally and copiously produced, contrary to the basic model assumptions.

The situation is somewhat paradoxically: the emission mechanism at the origin of the burst radiation must be very efficient, and the synchrotron and inverse Compton mechanisms by relativistic particles are indeed very efficient radiation processes, but just because of the very rapid cooling their predicted spectrum is too steep.

One is therefore forced to look for alternatives. In the dense photon environment of the internal shock scenario a

highly efficient viable alternative radiation mechanism may be Comptonization by a quasi-thermal particle distribution, as proposed by Ghisellini & Celotti (1999) (see also Thompson 1994; Liang 1997; Stern 1999; Liang et al. 1999). In this model, the conversion of bulk kinetic into random energy may still be due to shell–shell collisions, but the typical energy of the radiating particle is sub-relativistic, being fixed by the balance between the acceleration and the cooling processes. There is still equipartition between magnetic field, leptons and protons energies, but in a time integrated sense: all leptons are accelerated up to small energies, but for the entire duration of the shell–shell collision. Electron–positron pairs can be produced, and may even be the key ingredient to lock the particle energies in the observed range. These sub- or mildly relativistic particles would then emit self-absorbed cyclo-synchrotron photons and a Comptonization spectrum with a typical slope $F_\nu \propto \nu^0$ plus a Wien peak located where photon and particle energies are equal. The predicted spectral and temporal evolutions from this model are under investigation.

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